

Neutrino masses and Z' physics

V.V.Kiselev

Institute for High Energy Physics
142284, Protvino, Moscow Region, Russia.
E-mail: kiselev@mx.ihep.su

Abstract

The characteristic scale of neutral current, provided by an extension of Standard Model with a local group over the right-handed fermions, determines the smallness of neutrino masses of Dirac kind. The experimental observation of neutrino oscillations imposes the stringent limit on the Z' physics appearance at low energies.

1 Introduction

The recent clear observation of neutrino oscillations [1] indicates the non-zero masses for the neutrinos [2], which are generally considered to be extremely small due to the following mechanisms. The common form of neutrino mass matrix is expressed through three terms:

1. the Majorana mass m_T for the left-handed neutrino, that appears as the triplet weak isospin contribution,
2. the Dirac mass term m_D , involving the interaction with the sterile right-handed neutrino ν_R ,
3. the Majorana mass m_S for the singlet ν_R over the weak interaction.

Then the Lagrangian part, determining the neutrino mass, is equal to

$$L_m = (\bar{\nu}_R, \bar{\nu}_L) \begin{pmatrix} m_S & m_D \\ m_D & m_T \end{pmatrix} (1 + C) \begin{pmatrix} \nu_R \\ \nu_L \end{pmatrix} + \text{h.c.}, \quad (1)$$

where C denotes the charge conjugation. The following possibilities are generally discussed.

The first is the absence of sterile ν_R , that means there is the only contribution due to the Majorana term m_T . If it is caused by the vacuum expectation v of standard higgs, being the isodoublet, then the triplet mass is the square of v . Hence, $m_T \sim v^2/M$, where M appears from an extension of SM, and it is of large scale.

Second, the singlet contribution is determined by the physics beyond the SM, so that $m_S \gg m_D \gg m_T$, where m_D is usually taken in the range, corresponding to the mass scale for the charged fermions, $m_D \sim v$. Then the see-saw mechanism [3] leads to two species of neutrinos with the small and large masses, $m_1 \sim m_D^2/m_S$ and $m_2 \sim m_S$, correspondingly.

These scenarios (just beyond a super-physics) generally exhaust the natural explanations for the smallness of neutrino masses. Anyway, the experimental data put $m_\nu > 0.01$ eV, which means $M \sim m_S \sim 10^{15}$ GeV.

Note, that we do not know how the small Dirac masses can be naturally explained with no involvement of large sterile Majorana mass.

In the present paper we offer the scheme, wherein the neutrinos have zero masses in the Standard Model and acquire the smallness after an extension of SM to include the right-handed neutral currents. We suggest some non-trivial vacuum correlators, which determine the mass scales in connection to the gauge charges of fermions. Thus, it happens that the large scale M is related to the mass of Z' .

2 Mass generation

The second order contribution of SM to the effective action contains the neutral current term of the form

$$S_{2m} = \int dxdy \frac{e}{\cos \theta \sin \theta} (T_3 - Q \sin^2 \theta) [\bar{L}_L(x) Z_\mu(x) \gamma^\mu L_L(x)] \cdot Qe \tan \theta [\bar{L}_R(y) Z_\nu(y) \gamma^\nu L_R(y)], \quad (2)$$

where we introduce the notations L_L for the left-handed doublets and L_R for the right-handed singlets. Further, suggest the non-trivial vacuum correlators with the characteristic distance $r \sim 1/v$

$$\begin{aligned} \langle 0 | Z_\mu(x) \gamma^\mu L_L(x) \bar{L}_R(y) Z_\nu(y) \gamma^\nu | 0 \rangle &= \frac{\delta(x-y)}{v^4} \langle 0 | Z_\mu(x) \gamma^\mu L_L(x) \bar{L}_R(x) Z_\nu(x) \gamma^\nu | 0 \rangle \\ &\sim \delta(x-y) v, \end{aligned} \quad (3)$$

where we suppose that the scales of expectations for ZZ and $L_L \bar{L}_R$ equal v^2 and v^3 , respectively. Therefore, the fermion masses of Dirac kind are determined by the action

$$S_{fm} \sim \int dx \bar{L}_L(x) L_R(x) \cdot v \cdot \frac{e^2}{\cos^2 \theta} (T_3 - Q \sin^2 \theta) Q + \text{h.c.} \quad (4)$$

From (4) we deduce that the coupling of vacuum expectations, causing the Dirac masses, is determined by the charges of fermions, so that, say, for the neutrino the electric charge equal to zero results in the massless, which, thus, looks quite natural in the SM with the suggested mechanism for the mass generation.

Sure, we could introduce the local source, i.e. the Higgs field, for the vacuum expectation considered in the model above and make the Legendre transformation to substitute the field for the condensates. Then, we believe, the action would take the most usual form¹.

¹In this way we have, say, to consider the initial bi-local bare J -action

$$S_{ib} = \int dxdy J(x, y) [\bar{L}_R(x) Z_\mu(x) \gamma^\mu Z_\nu(y) \gamma^\nu L_L(y)] - \int dx \phi(x) J(x, x) + \text{h.c.}$$

To the bare order the equation of motion for the bi-local field results in the straightforward substitution of local field ϕ , as it stands in the above consideration for the correlators, developing the vacuum expectation values. After the analysis of divergences in the J -dependent Green functions, the corresponding counter terms must be added to the action. Then the J -source can be integrated out, that believes to result in the ϕ -higgs action, containing the couplings to fermions as well as the suitable potential to develop the spontaneous breaking of symmetry.

2.1 Neutrino masses from right-handed extension

Suppose that there is the additional $SU(2)_R$ local symmetry spontaneously broken at a scale v_R (for the sake of explicitness we put $v = v_L$ and $v_R^2 = b v_L^2 \gg v_L^2$). To minimize possible virtual corrections at low energies precisely studied up to the LEP measurements and to reproduce the mass relations between the standard gauge bosons, we have to introduce the additional Higgs field, being the doublet over $SU(2)_R$, which possesses zero charges over $U(1) \otimes SU(2)_L$ of the SM. The important challenge is that the standard higgs has to be extended to the field, belonging to the $(\frac{1}{2}, \frac{1}{2})$ representation of $SU(2)_L \otimes SU(2)_R$. The reason is the desirable renormalizability of the field theory with the spontaneously broken symmetry. Indeed, the fermion mass term $\bar{\psi}_R \psi_L$ certainly is of $(\frac{1}{2}, \frac{1}{2})$, so that the corresponding higgs developing the VEV must have the same quantum numbers over the local group².

Next, the non-zero VEVs of neutral Higgs fields result in the massless of the photon and lead to the massive neutral currents. The mass matrix for the local gauge fields B , Z_L and Z_R is determined by the form

$$M_{gauge}^2 = \frac{1}{4} v_L^2 \begin{pmatrix} g^2 & -gg_L & gg_R \\ -gg_L & g_L^2 & -g_L g_R \\ gg_R & -g_L g_R & g_R^2(1+b) \end{pmatrix}, \quad (5)$$

where g , g_L and g_R denote the gauge fields couplings for the $U(1) \otimes SU(2)_L \otimes SU(2)_R$ group.

Then the eigenvalues of the mass matrix are equal to

$$\begin{aligned} m_A^2 &= 0, \\ m_Z^2 &= \frac{1}{8} v_L^2 \left(\sum_i g_i^2 + b g_R^2 - \sqrt{(\sum_i g_i^2)^2 - 2b g_R^2 (g^2 + g_L^2 - g_R^2) + b^2 g_R^4} \right), \\ m_{Z'}^2 &= \frac{1}{8} v_L^2 \left(\sum_i g_i^2 + b g_R^2 + \sqrt{(\sum_i g_i^2)^2 - 2b g_R^2 (g^2 + g_L^2 - g_R^2) + b^2 g_R^4} \right), \end{aligned} \quad (6)$$

which in the limit of infinitely large b tend to the following relations:

$$\begin{aligned} m_A^2 &= 0, \\ m_Z^2 &\approx \frac{1}{4} v_L^2 (g^2 + g_L^2) \left(1 - \frac{1}{b}\right) \approx \frac{1}{4} v_L^2 (g^2 + g_L^2), \\ m_{Z'}^2 &\approx \frac{1}{4} [(v_R^2 + v_L^2) g_R^2 + v_L^2 (g^2 + g_L^2) \frac{1}{b}] \approx \frac{1}{4} v_R^2 g_R^2. \end{aligned} \quad (7)$$

Furthermore, it is quite evident to derive that the masses of charged gauged bosons are given by

$$\begin{aligned} m_{W_L}^2 &= \frac{1}{4} v_L^2 g_L^2, \\ m_{W_R}^2 &= \frac{1}{4} [v_R^2 g_R^2 + v_L^2 g_L^2]. \end{aligned}$$

²For the subject under consideration it is not so significant that the extension is $SU(2)_R$. It can be, say, $U(1)_R$. In the current discussion we can suppose that the right-handed doublets are $\begin{pmatrix} e \\ d \end{pmatrix}$ and $\begin{pmatrix} \nu \\ u \end{pmatrix}$. Then the charged $SU(2)_R$ -like gauge bosons are the vector leptoquarks, that can be the reason for the very different physics in the left- and right-handed sectors.

Thus, we see that up to small corrections the Z boson mass reproduces the value of standard boson as it is connected to the W mass.

The matrix S , transforming the gauge fields to the mass eigen-states, has the form

$$\begin{pmatrix} A \\ Z \\ Z' \end{pmatrix} = \begin{pmatrix} \cos \theta & \sin \theta & 0 \\ -\sin \theta & \cos \theta & \frac{1}{\beta} \\ \frac{1}{\beta} \sin \theta & -\frac{1}{\beta} \cos \theta & 1 \end{pmatrix} \begin{pmatrix} B \\ Z_L \\ Z_R \end{pmatrix}, \quad (8)$$

with the accuracy up to $O(\frac{1}{\beta^2})$, where $\beta = bg_R/\sqrt{g^2 + g_L^2}$, and θ is the standard angle by Weinberg. In the approximation under consideration we can see that S has the orthogonal form, and the transposition results in the inverse matrix,

$$\begin{pmatrix} B \\ Z_L \\ Z_R \end{pmatrix} = \begin{pmatrix} \cos \theta & -\sin \theta & \frac{1}{\beta} \sin \theta \\ \sin \theta & \cos \theta & -\frac{1}{\beta} \cos \theta \\ 0 & \frac{1}{\beta} & 1 \end{pmatrix} \begin{pmatrix} A \\ Z \\ Z' \end{pmatrix}, \quad (9)$$

so that the admixture of Z' in Z_L is determined by the ratio of v_L^2/v_R^2 .

The vertices of massive eigen-states are determined by the following relations:

1. The photon couples to the electric charge

$$Q = \frac{Y^L}{2} + T_3^L = \frac{Y^R}{2}.$$

2. The left-handed fermions have the standard couplings to Z

$$\frac{e}{\cos \theta \sin \theta} (T_3^L - Q \sin^2 \theta).$$

3. The right-handed fermions acquire the correction to the charge with Z

$$-Qe \tan \theta + T_3^R \frac{g_R}{\beta}.$$

4. The Z' vertex to the left-handed fermions is proportional to that of Z one, so that

$$-\frac{e}{\beta \cos \theta \sin \theta} (T_3^L - Q \sin^2 \theta),$$

and the suppression is due to the smallness of v_L^2/v_R^2 .

5. Z' has the charge

$$T_3^R g_R + Qe \frac{\tan \theta}{\beta},$$

to the right-handed fermions.

So, some anomalous couplings are introduced due to the Z' physics.

Now suggesting the non-trivial vacuum correlations at the distances of $r \sim 1/v_R$ we find that the neutrinos acquire the non-zero masses due to the admixture of Z' in Z_L and its dominance in Z_R . So, the action

$$S'_{2m} = - \int dx dy \frac{e}{\beta \cos \theta \sin \theta} (T_3^L - Q \sin^2 \theta) [\bar{L}_L(x) Z'_\mu(x) \gamma^\mu L_L(x)] \cdot \\ \left(T_3^R g_R + Q e \frac{\tan \theta}{\beta} \right) [\bar{L}_R(y) Z'_\nu(y) \gamma^\nu L_R(y)], \quad (10)$$

transforms to

$$S'_{fm} \sim \int dx \bar{L}_L(x) L_R(x) \cdot \frac{v_L^2}{v_R} \cdot \frac{e g_R}{\beta \cos \theta \sin \theta} (T_3^L - Q \sin^2 \theta) T_3^R + \text{h.c.} \quad (11)$$

if

$$\langle 0 | Z'_\mu(x) \gamma^\mu L_L(x) \bar{L}_R(y) Z'_\nu(y) \gamma^\nu | 0 \rangle = \frac{\delta(x-y)}{v_R^4} \langle 0 | Z'_\mu(x) \gamma^\mu L_L(x) \bar{L}_R(x) Z'_\nu(x) \gamma^\nu | 0 \rangle \\ \sim \delta(x-y) v_R, \quad (12)$$

where we suppose that the scales of expectations in these correlations for $Z'Z'$ and $L_L \bar{L}_R$ equal v_R^2 and v_R^3 , respectively.

Note, that, first, other correlations result in less contributions to the masses, as those are suppressed by powers of v_L/v_R . Second, the corrections to the masses of electrically charged fermions seem to be suppressed in the same manner.

Thus, we see that due to the extension of model to the right-handed local group the neutrinos have the Dirac mass of the order of $m_\nu \sim v_L^2/v_R$, where-from we extract $v_R \sim 10^{15}$ GeV. If v_R is so large, the current experimental bounds on the anomalous couplings of gauge bosons and the appearance of Z' [4] are far away from what is expected from the small neutrino masses. The other possibility is to assume the existence of additional sterile neutrino mass m_S , which activates the see-saw mechanism, but in the model, where the Dirac mass can be essentially reduced from 200 GeV by several orders of magnitude due to the suppression v_L/v_R .

3 Hierarchy of scales and GUT

The arrangement of vacuum expectation values for the spontaneous breaking of the local gauge symmetries can be reasonably related to the following qualitative peculiarities, belonging to the corresponding invariant actions. So, we observe:

- The abelian $U(1)$ -field, coupled in the vector-like way to the fermions of both chiralities, does not appear in the spontaneously breaking phase.
- The non-abelian $SU(3)$ field, possessing the asymptotic freedom, is coupled, again, in the vector-like way to the fermions, and it does not acquire the spontaneous breaking of symmetry, too. However, the back-wise face of asymptotic freedom is the confinement.

- The non-abelian $SU(2)$ -field, coupled to the chiral fermions, exposes the spontaneous breaking.

The presence of asymptotic freedom for the latter symmetry depends on the set of matter fields. In what follows, we exploit the situation, when $SU(2)$ is asymptotically free.

Let us offer the following picture. The structure of vector-like gauge symmetries, i.e. the form of effective potential, preserves them from the developing of non-trivial vacuum correlations, determining the spontaneous breaking of invariances. The non-abelian theory with chiral fermions does possess the effective action, where the vacuum correlations appear, if the coupling constant is greater than a critical value $\hat{\alpha}$. Then in GUT [5] with $\alpha_{\text{GUT}} < \hat{\alpha}$, the chiral $SU(2)$ will develop the symmetry breaking VEV at a low scale M_2 , where its coupling $\alpha_2(M_2)$ will reach the critical value. That can be the reason for the very different values of M_{GUT} and M_2 . Note, that in this approach we know the value of critical constant $\hat{\alpha}$, since it is well measured in the weak interactions, so that $\hat{\alpha} \approx 1/30$.

As for the neutrino mass generation described above, we are ready to conclude, that the difference between the renormalization group properties, determining the coupling running, for the left- and right-handed symmetries will result in the hierarchy of scales for their characteristic VEVs.

To be more concrete, consider the gauge symmetry for the right-handed fermions, embedded to the following $SU(2)_R$ -like doublets:

$$\begin{pmatrix} e \\ d \end{pmatrix}_R \quad \begin{pmatrix} \nu \\ u \end{pmatrix}_R \quad (13)$$

The essential difference from the usual $SU(2)$ is that the charged vector bosons, possessing the fractional electric charge $\pm\frac{2}{3}$, are the color triplet 3_c and anti-triplet $\bar{3}_c$, appropriately. The corresponding generalized derivative, acting on the fermions, has the form

$$iD^\mu = i\partial^\mu - \frac{g_R}{2}(\tau_3 Z_R^\mu + \sqrt{2}\tau_+ W_{-2/3}^\mu + \sqrt{2}\tau_- W_{+2/3}^\mu), \quad (14)$$

where $\tau_{3,\pm}$ are the Pauli matrices, and the superscript i runs over the color anti-triplet and the subscript i does the color triplet.

It is quite clear, that, including the quark colors, the "doublets" in (13) could represent the $SU(4)_R$ fundamental multiplet, if the couplings of g_R for the right-handed fermions and g_3 in $SU(3)$ would be equal each to other. If we "switch off" the color indexes from the group transformations, the corresponding invariance G_R is $SU(4)/SU(3)$ on the right-handed fermions.

As we have mentioned, the G_R symmetry is very similar to the famous $SU(2)$. The straightforward consideration leads to that we can reproduce the one-loop calculations for the running of g_R from the corresponding evaluation for $SU(2)$, if we substitute for $C_F = \frac{(N^2-1)}{2N} = \frac{3}{4}$ at $N = 2$ by $\tilde{C}_F = \frac{1+2N_c}{4} = \frac{7}{4}$, where $N_c = 3$ is the number of colors, and for $C_A = N = 2$ by $\tilde{C}_A = 2N_c = 6$.

The running of coupling constants in the $SU(N)$ field theory is given by the expression

$$\frac{1}{\alpha_N(M)} = \frac{1}{\alpha_N(M_0)} + \frac{b_N}{2\pi} \ln \frac{M}{M_0}, \quad (15)$$

where b_N depends on the set of fields. So,

$$b_N = \frac{11}{3}N - \frac{1}{3}n_f - \frac{1}{6}n_s, \quad (16)$$

where n_f is the number of chiral fermions, n_s is the number of fundamental scalar multiplets. Therefore, for the b -coefficients of $G_R \otimes SU(2)_L \otimes SU(3)$ we get

$$\begin{aligned} b_R &= 22 - \frac{2}{3}n_g - \frac{1}{6}n_{s(R)}, \\ b_L &= \frac{22}{3} - \frac{4}{3}n_g - \frac{1}{6}n_{s(L)}, \end{aligned} \quad (17)$$

$$b_3 = 11 - \frac{4}{3}n_g - \frac{1}{6}n_{s(c)}, \quad (18)$$

where n_g is the number of fermion generations, $n_{s(R,L,c)}$ are the numbers of corresponding scalars.

Next, the $U(1)$ coupling constant, normalized as $\alpha_1 = \frac{5}{3}\alpha_Y$, where Y is the weak hyper-charge, has the coefficient b_1 equal to

$$b_1 = -\frac{4}{3}n_g - \frac{2}{15}n_{s(Y)}, \quad (19)$$

where we have taken into account the hyper-charge of additional weak doublet due to the extension of standard higgs by G_R :

$$\begin{pmatrix} h_+ \\ h_0 \end{pmatrix} \rightarrow \begin{pmatrix} h_+ & h_{+1/3} \\ h_0 & h_{-2/3} \end{pmatrix}, \quad (20)$$

where the fractionally charged higgses are the color anti-triplets with the hyper-charge $Y = -1/3$. So, $n_{s(Y)}$ denotes the number of standard higgses.

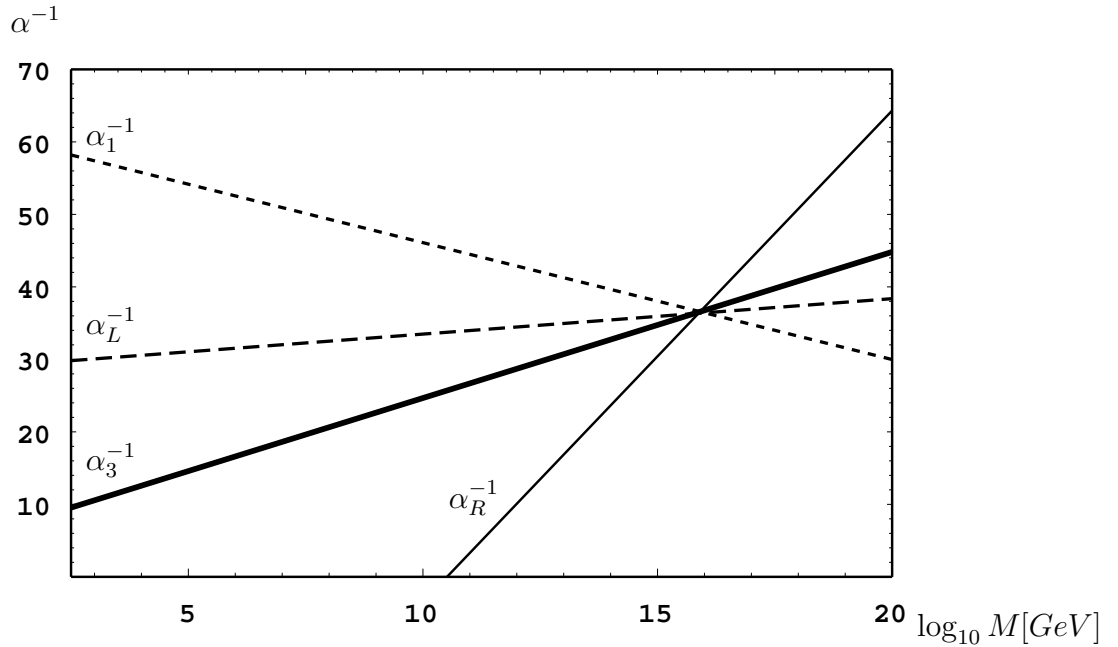


Figure 1: The unification of couplings in GUT with no SUSY.

Let us discuss the scalar field set, suitable for the problem. First, the number of left-symmetric higgses includes $1 + 3$ doublets. In realistic models for the generation replication and their mixing,

the number of standard higgses usually repeats the number of generations [6]. So, we put

$$n_{s(L)} = 4n_g, \quad n_{s(Y)} = n_g. \quad (21)$$

Second, the right-symmetric scalars are those of extensions for the standard higgs and that of $SU(2)_L$ singlet to separately break the G_R -symmetry. We put

$$n_{s(R)} = 3n_g, \quad n_{s(c)} = 3n_g. \quad (22)$$

Further, we can look at the evolution of couplings to large scales as it is shown in Fig.1, and draw the conclusion on the plausible unification of symmetries at $M_{\text{GUT}} \sim 8 \cdot 10^{15}$ GeV. We present this picture for the illustration of other feature: the α_R coupling reaches the critical region, $\alpha \simeq \hat{\alpha} = 1/30$, at the scales, which are only one or two orders of magnitude less than the GUT energy because of the appropriate properties in the renormalization group. Of course, the numerical estimate qualitatively depends on $\hat{\alpha}$, which can vary over the structure of right-handed symmetry. Another note concerns the extra-higgses, which, according to the evolution performed, are much lighter than M_{GUT} .

As for the model under discussion, we could add only that, obviously, there is the right-handed symmetry at M_{GUT} : $SU(4)_R$ with the violation in the way $SU(4)_R \otimes SU(3)_L \rightarrow G_R \otimes SU(3)$.

To complete, we have to emphasize that the correlators in (10)-(12) belong to the $(\frac{1}{2}, \frac{1}{2})$ representation over $G_R \otimes SU(2)_L$, and, hence, contribute, a little bit, to VEV of the extended higgs, and not to the G_R -doublet, developing the v_R scale itself.

4 Conclusion

We have shown how the model of mass generation can be constructed on the basis of higgs mechanism, wherein the scalar isodoublet field is related to the vacuum correlations and fermion charges, so that

- the neutrino is massless in the Standard Model because of its zero electric charge,
- the smallness of neutrino masses can be caused by the hierarchy of the correlation scales for the spontaneous breaking of the standard local symmetry and the right-handed extension, $v_L \ll v_R$, which leads to the Dirac kind of mass,
- the unification of coupling constants makes v_R to be only one or two orders of magnitude less than M_{GUT} ,
- we could reduce the scale of sterile Majorana mass by involving both the see-saw mechanism and suppressed Dirac terms due to the above approach.

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